

Electroweak baryogenesis and quantum corrections to the triple Higgs boson coupling

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Abstract

A phenomenological consequence of electroweak baryogenesis is studied in connection with the Higgs physics. In a two Higgs double model, the first order phase transition can be strong enough to allow the electroweak baryogenesis due to the effect of extra Higgs bosons. We investigate the quantum correction to the triple coupling of the lightest Higgs boson in such a scenario, and find that the condition of the strong first order phase transition necessarily leads to the deviation of at least 10% from the standard model prediction. Such magnitude of the deviation can be identified in future e^+e^- linear collider experiments. We also discuss the minimal supersymmetric standard model with the light stop scenario, and point out that a sizable correction appears in the triple coupling for successful electroweak baryogenesis.

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One of the most challenging problems in particle physics and cosmology is to explain the baryon asymmetry of the Universe, $n_B/s \sim 10^{-10}$ [1], where n_B is the difference between the number density of baryons and that of anti-baryons, and s is the entropy density. Three conditions are required for generation of baryon asymmetry starting from a baryon-symmetric initial state [2]: (1) baryon number nonconservation, (2) C and CP violation and (3) departure from thermal equilibrium. These conditions can be in principle satisfied at the electroweak phase transition, because baryon number violating processes are efficient at high temperature [3]. In the minimal Standard Model (SM), however, it turns out to be difficult to generate sufficient baryon asymmetry observed today, since the CP violating phase from Cabibbo-Kobayashi-Maskawa matrix is too small and the strength of the phase transition is too weak for the current Higgs boson mass bound ($m_h \gtrsim 114$ GeV [4]). To realize the baryogenesis at the electroweak phase transition we must extend the minimal SM Higgs sector. Various scenarios have been considered for electroweak baryogenesis [5].

In this Letter, we consider the electroweak baryogenesis in the Two Higgs Doublet Model (THDM) [6, 7, 8, 9]. In particular, we focus on the connection between cosmology and collider physics. We study critical temperature and the order of phase transition by the use of the one-loop effective potential at finite temperature. It is found that the phase transition can be strongly first order enough for baryogenesis to take place owing to the non-decoupling effect of the heavy Higgs bosons. We also calculate the radiative correction to the triple coupling (hhh) of the lightest Higgs boson. We find that the deviation of the hhh coupling from the SM value is larger than about 10 percent for $m_h = 120$ GeV when the necessary condition for the electroweak baryogenesis is satisfied. Since the experimental accuracy of the hhh coupling determination is expected to be $\mathcal{O}(10)\%$ from the double Higgs production process at a future e^+e^- Linear Collider (LC) [10], the self-coupling measurement can provide an important test of the electroweak baryogenesis scenario ¹. We then extend our consideration to the case of minimal supersymmetric standard model (MSSM).

In order to create the large enough baryon number asymmetry during the electroweak phase transition, we need two conditions on the Higgs potential. First, the phase transition must be strongly first order to avoid the washout of the generated baryon number density by the sphaleron process after the phase transition [12]. Secondly, there should be CP violating

¹ The preliminary result is reported in Ref. [11].

phases in the Higgs potential at finite temperature, which are necessary to generate asymmetry of the fermion and anti-fermion around the bubble wall during the phase transition. The CP violating phases from the Higgs sector originate either from explicit or spontaneous CP violation. Therefore, connection between the CP phases at the phase transition and those in the zero-temperature effective potential strongly depends on the scenario for the electroweak baryogenesis. On the other hand, the strong first-order phase transition is a necessary requirement independent of the detail of the mechanism for baryogenesis. Hence, we here focus on the first condition for the finite temperature effective potential [13].

In the THDM the tree-level Higgs potential is given by

$$V_{\text{tree}} = m_1^2 |\Phi_1|^2 + m_2^2 |\Phi_2|^2 - (m_3^2 \Phi_1^\dagger \Phi_2 + \text{h.c.}) + \frac{\lambda_1}{2} |\Phi_1|^4 + \frac{\lambda_2}{2} |\Phi_2|^4 + \lambda_3 |\Phi_1|^2 |\Phi_2|^2 + \lambda_4 |\Phi_1^\dagger \Phi_2|^2 + \left[\frac{\lambda_5}{2} (\Phi_1^\dagger \Phi_2)^2 + \text{h.c.} \right], \quad (1)$$

where we imposed the softly-broken discrete symmetry as $\Phi_1 \rightarrow \Phi_1$, $\Phi_2 \rightarrow -\Phi_2$ to suppress flavor changing neutral current processes [14]². Here m_3^2 and λ_5 are generally complex. There is one physical phase which cannot be removed by rephasing of the doublet fields, and this becomes the source of explicit CP violation. Although this phase may be important for generation of baryon asymmetry, we neglect it in the following discussion on the electroweak phase transition³. In the limit of CP invariance there are five physical states; CP-even (h, H), CP-odd (A) and charged Higgs bosons (H^\pm). Here, h is lighter than H . We introduce mixing angles α and β . The angle α is defined as the mixing angle between CP-even Higgs bosons, and $\tan \beta (= \langle \Phi_2^0 \rangle / \langle \Phi_1^0 \rangle)$ is the ratio of the vacuum expectation values (VEVs) of the two Higgs fields. These VEVs satisfy $v = \sqrt{2} \sqrt{\langle \Phi_1^0 \rangle^2 + \langle \Phi_2^0 \rangle^2} \simeq 246$ GeV. In addition, we introduce $M^2 (= m_3^2 / \sin \beta \cos \beta)$, which characterizes the soft breaking scale of the discrete symmetry. There are two cases concerning the decoupling property of the additional heavy Higgs bosons. When $M^2 \gg v^2$, all the masses of the heavy Higgs bosons are approximately given by M , and the properties of the lightest Higgs boson become similar to those of the SM Higgs boson. In this case the loop effect due to the heavy Higgs bosons is suppressed by the factor of $1/M^2$. On the contrary, in the opposite limit where all the heavy Higgs bosons

² According to the types of Yukawa interaction we can categorize the THDM into so called Model I and Model II. In Model I all fermions receive their masses from only one Higgs field, whereas in Model II the up-type quarks receive their masses from Φ_2 and the down-type quarks and leptons do from Φ_1 . The following discussion does not depend on the choice of Model I or Model II.

³ This is justified as long as the small phase approximation is valid to solve the bubble wall equations [8, 9].

receive their masses from the VEVs, the loop effect due to the heavy Higgs bosons does not necessarily decouple in the large mass limit.

In the following, we assume $m_1 = m_2$ and $\lambda_1 = \lambda_2$ in order to simplify the analysis of the phase transition. In this case, the field direction relevant to the electroweak phase transition is reduced to

$$\langle \Phi_1 \rangle = \langle \Phi_2 \rangle = \begin{pmatrix} 0 \\ \frac{1}{2}\varphi \end{pmatrix}. \quad (2)$$

This case corresponds to $\sin(\alpha - \beta) = -1$ and $\tan\beta = 1$. This parameter choice has been proposed for a successful scenario of electroweak baryogenesis [7, 8] to avoid the complexity with two stage phase transition [15].

The quantum correction of the extra Higgs bosons to the hhh coupling is calculated in the THDM using both effective potential and diagrammatic methods in Refs. [16, 17]. The effective potential is given by $V_{\text{eff}}[\varphi] = V_{\text{tree}}[\varphi] + \Delta V[\varphi]$, where $\Delta V[\varphi]$ is expressed as

$$\Delta V[\varphi] = \frac{1}{64\pi^2} \sum_f N_{c_f} N_{s_f} (-1)^{2s_f} (M_f[\varphi])^4 \left\{ \ln \frac{(M_f[\varphi])^2}{Q^2} - \frac{3}{2} \right\}. \quad (3)$$

Here N_{c_f} is the color number, s_f (N_{s_f}) is the spin (degree of freedom) of the field f in the loop, $M_f[\varphi]$ is the field dependent mass of f , and Q is a renormalization scale. In Ref. [16], we found that the deviation of the hhh coupling constant from SM value can be $\mathcal{O}(100)\%$ due to the non-decoupling effect of the heavy Higgs bosons. When the Higgs couplings to gauge bosons are SM-like at the tree-level ($\sin(\alpha - \beta) \simeq -1$), the leading loop contributions of the heavy Higgs bosons and the top quark are extracted as

$$\lambda_{hhh}^{\text{eff}}(\text{THDM}) \simeq \frac{3m_h^2}{v} \left[1 + \frac{m_H^4}{12\pi^2 m_h^2 v^2} \left(1 - \frac{M^2}{m_H^2} \right)^3 + \frac{m_A^4}{12\pi^2 m_h^2 v^2} \left(1 - \frac{M^2}{m_A^2} \right)^3 + \frac{m_{H^\pm}^4}{6\pi^2 m_h^2 v^2} \left(1 - \frac{M^2}{m_{H^\pm}^2} \right)^3 - \frac{m_t^4}{\pi^2 m_h^2 v^2} \right]. \quad (4)$$

It is easily seen that the effects of the heavy Higgs boson loops are enhanced by m_Φ^4 ($\Phi = H, A, H^\pm$) when M^2 is zero. These effects do not decouple in the large mass limit $m_\Phi \rightarrow \infty$ and yields the large deviation of hhh coupling from the SM prediction.

In order to study the electroweak phase transition, we consider the effective potential at finite temperature:

$$V_{\text{eff}}(\varphi, T) = V_{\text{tree}}(\varphi) + \Delta V(\varphi) + \Delta V^{(T)}(\varphi, T), \quad (5)$$

where the finite temperature contribution is expressed as

$$\Delta V^{(T)}(\varphi, T) = \frac{T^4}{2\pi^2} \left[\sum_{i=\text{bosons}} n_i I_B(a_i^2) + n_t I_F(a_t^2) \right], \quad (6)$$

with

$$I_{B,F}(a^2) = \int_0^\infty dx \, x^2 \log(1 \mp e^{-\sqrt{x^2+a^2}}). \quad (7)$$

In the above expressions, n_i is the degree of freedom of the particle i : $n_W = 6$, $n_Z = 3$, $n_t = -12$, $n_h = n_H = n_A = 1$, $n_{H^\pm} = 2$, $m_i(\varphi)$ is the field dependent mass and $a_i(\varphi) = m_i(\varphi)/T$.

Before going to the numerical evaluation, we discuss qualitative features of the phase transition by using the high temperature expansion. For high temperature $T \gg m_i$, we may expand the integrals $I_{B,F}$ analytically [18]. When $m_\Phi^2 \gg m_h^2, M^2$ ($\Phi = H, A, H^\pm$), the field dependent masses of the heavy Higgs bosons can be written as $m_\Phi^2(\varphi) \simeq m_\Phi^2 \varphi^2/v^2$. In the high temperature approximation, the effective potential at finite temperature of Eq. (6) is expanded as [19]

$$V_{\text{eff}}(\varphi, T) \simeq D(T^2 - T_0^2)\varphi^2 - ET|\varphi|^3 + \frac{\lambda_T}{4}\varphi^4 + \dots, \quad (8)$$

where

$$D = \frac{1}{24v^2}(6m_W^2 + 3m_Z^2 + 6m_t^2 + m_H^2 + m_A^2 + 2m_{H^\pm}^2), \quad (9)$$

$$T_0^2 = \left\{ \frac{1}{4}m_h^2 - \frac{1}{32\pi^2v^2}(6m_W^4 + 3m_Z^4 - 12m_t^4 + m_H^4 + m_A^4 + 2m_{H^\pm}^4) \right\} / D, \quad (10)$$

$$E = \frac{1}{12\pi v^3}(6m_W^3 + 3m_Z^3 + m_H^3 + m_A^3 + 2m_{H^\pm}^3), \quad (11)$$

$$\lambda_T = \frac{m_h^2}{2v^2} \left[1 - \frac{1}{8\pi^2v^2m_h^2} \left\{ 6m_W^4 \log \frac{m_W^2}{\alpha_B T^2} + 3m_Z^4 \log \frac{m_Z^2}{\alpha_B T^2} - 12m_t^4 \log \frac{m_t^2}{\alpha_F T^2} \right. \right. \\ \left. \left. + m_H^4 \log \frac{m_H^2}{\alpha_B T^2} + m_A^4 \log \frac{m_A^2}{\alpha_B T^2} + 2m_{H^\pm}^4 \log \frac{m_{H^\pm}^2}{\alpha_B T^2} \right\} \right], \quad (12)$$

where $\log \alpha_B = 2 \log 4\pi - 2\gamma_E$, $\log \alpha_F = 2 \log \pi - 2\gamma_E$ and γ_E is the Euler constant. The first order phase transition becomes possible due to appearance of the cubic term which originates from the boson loops at finite temperature. Unlike the SM, the coefficient of the cubic term E can receive a large contribution from the extra Higgs bosons. At the critical temperature T_c , the potential has two degenerate minima at $\varphi = 0$ and $\varphi_c = 2ET_c/\lambda_{T_c}$, where λ_{T_c} is the quartic coupling constant at T_c . In order not to wash out the created baryon number density after the electroweak phase transition, we have to require that the

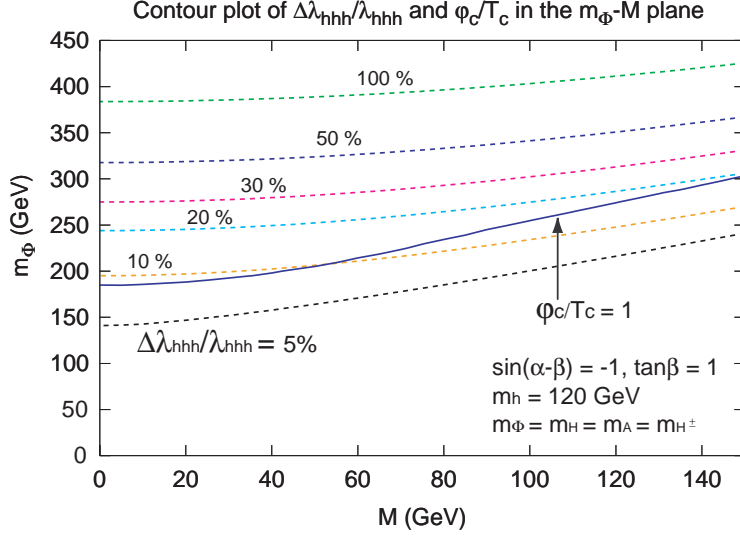


FIG. 1: The straight line stands for the critical line which satisfied the condition, $\varphi_c/T_c = 1$. The dashed lines are the deviation of hhh coupling from the SM value, where $\Delta\lambda_{hhh}^{\text{THDM}} \equiv \lambda_{hhh}^{\text{eff}}(\text{THDM}) - \lambda_{hhh}^{\text{eff}}(\text{SM})$.

sphaleron process should be sufficiently suppressed. The most reliable condition has been obtained from the lattice simulation study [20]. It is expressed as

$$\frac{\varphi_c}{T_c} = \frac{2E}{\lambda_{T_c}} \gtrsim 1. \quad (13)$$

For $m_h = 120$ GeV, this condition can be satisfied when the masses of the heavy Higgs bosons are above 200 GeV. We can see from Eq. (4) that the correction to the hhh coupling can be large in such a parameter region. Although the high temperature expansion gives a qualitative description of the phase transition, the approximation breaks down when the masses of the heavy Higgs bosons become larger than the critical temperature. We therefore evaluate the effective potential numerically and search the parameter space where the condition (13) is satisfied.

In Fig. 1, we show the parameter region where the necessary condition of the electroweak baryogenesis in Eq. (13) is satisfied in the m_Φ - M plane. We take $\sin(\alpha - \beta) = -1$, $\tan \beta = 1$ and $m_h = 120$ GeV. For the heavy Higgs boson mass, we assume $m_H = m_A = m_{H^\pm} (\equiv m_\Phi)$ to avoid the constraint on the ρ parameter from the LEP precision data [21]. In the numerical evaluation, we take into account the ring summation for the contribution of the Higgs bosons to the effective potential at finite temperature [18, 22]. For fixed values of m_Φ and M , we calculate the effective potential (6) varying the temperature T and determine the

critical temperature T_c of the first-order phase transition and the expectation value φ_c at T_c . We can see that the phase transition becomes strong enough for the baryogenesis when the masses of the heavy Higgs bosons are larger than about 200 GeV. For larger values of M , greater m_Φ are required to satisfy the condition (13). In this figure we also plot the contour of the magnitude of the deviation in the hhh coupling from the SM value. We define the deviation $\Delta\lambda_{hhh}^{\text{THDM}}/\lambda_{hhh}^{\text{eff}}(\text{SM})$ by $\Delta\lambda_{hhh}^{\text{THDM}} \equiv \lambda_{hhh}^{\text{eff}}(\text{THDM}) - \lambda_{hhh}^{\text{eff}}(\text{SM})$. We calculated the deviation at the one loop level in the on-shell scheme which gives a better approximation than the use of Eq. (4) [16, 17]. We can easily see that the magnitude of the deviation is significant ($\gtrsim 10\%$) in the parameter region where the electroweak baryogenesis is possible. Such magnitude of the deviation can be detected at a future LC experiment.

Next we discuss a scenario of electroweak baryogenesis in the MSSM. In this case, the strong first order phase transition can be induced by the loop effect of the light stop in the finite temperature effective potential [23, 24]. We examine the loop effect of the light stop on the hhh coupling in this scenario. In the following, we only consider the finite and zero temperature effective potentials in a simple approximation to understand the qualitative feature.

In order to realize the successful baryogenesis under the constraint from the ρ parameter and the lower bound of the lightest Higgs boson mass, we have to take the $M_Q^2 \gg M_U^2, m_t^2$, where M_Q and M_U are soft breaking masses associated with left and right stops, respectively. The eigenvalues of the field dependent mass matrix are expressed as

$$m_{\tilde{t}_1}^2(\varphi, \beta) = M_U^2 + D_R^2 + \frac{h_t^2 \sin^2 \beta}{2} \left(1 - \frac{|X_t|^2}{M_Q^2} \right) \varphi^2, \quad (14)$$

$$m_{\tilde{t}_2}^2(\varphi, \beta) = M_Q^2 + D_L^2 + \frac{h_t^2 \sin^2 \beta}{2} \left(1 + \frac{|X_t|^2}{M_Q^2} \right) \varphi^2 \simeq M_Q^2, \quad (15)$$

where $X_t = A_t + \mu \cot \beta$ with A_t and μ being the stop trilinear coupling and the higgsino mass term, respectively, and D_L and D_R are corresponding D-term contributions. In the high temperature expansion, the contribution to the coefficient of the cubic term in the finite temperature effective potential (8) is given by [23]

$$\Delta E_{\tilde{t}_1} \simeq + \frac{1}{6\pi} \frac{N_c m_t^3}{v^3} \left(1 - \frac{|X_t|^2}{M_Q^2} \right)^{3/2}, \quad (16)$$

for $M_U^2 + D_R^2 \sim 0$. It is known that electroweak baryogenesis is possible due to the stop contribution if the lightest Higgs mass m_h is close to the current experimental lower bound.

We also evaluate the one loop effect of the light stop to the hhh coupling, and extract the leading contribution as

$$\frac{\Delta\lambda_{hhh}(MSSM)}{\lambda_{hhh}(SM)} \simeq \frac{N_c m_t^4}{6\pi^2 v^2 m_h^2} \left(1 - \frac{|X_t|^2}{M_Q^2}\right)^3, \quad (17)$$

where m_h is the one-loop corrected mass of the lightest Higgs boson. Combining the above equations, we can derive

$$\frac{\Delta\lambda_{hhh}(MSSM)}{\lambda_{hhh}(SM)} \simeq \frac{6v^4}{m_t^2 m_h^2 N_c} (\Delta E_{t_1})^2. \quad (18)$$

From the condition (13), the deviation in the hhh coupling from the SM value is estimated to be larger than $\sim 6\%$. Although the above calculation needs to be improved, this example suggests that a large correction in the hhh coupling is a general feature of successful electroweak baryogenesis.

We have studied the electroweak phase transition by using the effective potential at finite temperature in the THDM and the MSSM. In the THDM, it is found that the phase transition can be strongly first order enough for successful electroweak baryogenesis due to the non-decoupling effects of the heavy Higgs boson loops at finite temperature. Such non-decoupling effects also affect the triple Higgs self-coupling constant of the lightest Higgs boson at the one-loop level. We found that if the baryogenesis occurs at the electroweak phase transition, the hhh coupling constant deviates from the SM prediction at least $\sim 10\%$ in the THDM. The measurement of the Higgs self-coupling therefore can provide an interesting interplay between cosmology and collider physics.

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